

Standard-smooth hybrid inflation

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(Dated: February 2, 2008)

We consider the extended supersymmetric Pati-Salam model which, for $\mu > 0$ and universal boundary conditions, succeeds to yield experimentally acceptable b -quark masses by moderately violating Yukawa unification. It is known that this model can lead to new shifted or new smooth hybrid inflation. We show that a successful two-stage inflationary scenario can be realized within this model based only on renormalizable superpotential interactions. The cosmological scales exit the horizon during the first stage of inflation, which is of the standard hybrid type and takes place along the trivial flat direction with the inflaton driven by radiative corrections. Spectral indices compatible with the recent data can be achieved in global supersymmetry or minimal supergravity by restricting the number of e-foldings of our present horizon during the first inflationary stage. The additional e-foldings needed for solving the horizon and flatness problems are naturally provided by the second stage of inflation, which occurs mainly along the built-in new smooth hybrid inflationary path appearing right after the destabilization of the trivial flat direction at its critical point. Monopoles are formed at the end of the first stage of inflation and are, subsequently, diluted by the second stage of inflation to become utterly negligible in the present universe for almost all (for all) the allowed values of the parameters in the case of global supersymmetry (minimal supergravity).

PACS numbers: 98.80.Cq

I. INTRODUCTION

It is well known [1] that the standard supersymmetric (SUSY) realization [2, 3] of hybrid inflation [4] in the context of grand unified theories (GUTs) leads, at the end of inflation, to a copious production of topological defects such as cosmic strings [5], magnetic monopoles [6], or domain walls [7] if these defects are predicted by the underlying symmetry breaking. In the case of magnetic monopoles or domain walls, this causes a cosmological catastrophe. The simplest GUT gauge group whose breaking to the standard model (SM) gauge group G_{SM} predicts the existence of topologically stable magnetic monopoles is the Pati-Salam (PS) group $G_{\text{PS}} = \text{SU}(4)_c \times \text{SU}(2)_L \times \text{SU}(2)_R$ [8]. (Note that the PS monopoles carry [9] two units of Dirac magnetic charge.) So, applying the standard realization of hybrid inflation within the SUSY PS GUT model, we encounter a cosmologically disastrous overproduction of magnetic monopoles at the end of inflation, where the GUT gauge symmetry G_{PS} breaks spontaneously to G_{SM} .

Possible ways out of this difficulty are provided by the shifted [10] or smooth [1, 11] variants of SUSY hybrid inflation, which, in their conventional realization, utilize non-renormalizable superpotential terms (for a review, see Ref. [12]). In these inflationary scenarios, the GUT gauge symmetry G_{PS} is broken to G_{SM} already during inflation and, thus, no magnetic monopoles are produced at the termination of inflation.

It has been shown [13, 14] that hybrid inflation of

both the shifted and smooth type can be implemented within an extended SUSY PS model without the need of non-renormalizable superpotential interactions. It is very interesting to note that this extended SUSY PS model was initially constructed [15] (see also Ref. [16]) for solving a very different problem. In SUSY models with exact Yukawa unification [17], such as the simplest SUSY PS model (see Ref. [18]), and universal boundary conditions, the b -quark mass comes out [19] unacceptably large for $\mu > 0$. Therefore, Yukawa unification must be moderately violated so that, for $\mu > 0$, the predicted b -quark mass resides within the experimentally allowed range even with universal boundary conditions. This requirement forces us to extend the superfield content of the SUSY PS model by including, among other superfields, an extra pair of $\text{SU}(4)_c$ non-singlet $\text{SU}(2)_L$ doublets, which naturally develop [20] subdominant vacuum expectation values (VEVs) and mix with the main electroweak doublets of the model leading to a moderate violation of Yukawa unification. It is quite remarkable that the resulting extended SUSY PS model can automatically and naturally lead [13, 14] to a new version of shifted and smooth hybrid inflation based solely on renormalizable superpotential terms. As in the conventional realization of shifted and smooth hybrid inflation, the GUT gauge group G_{PS} is broken to G_{SM} already during inflation in the new shifted [13] and new smooth [14] hybrid inflation scenario too and monopole production at the end of inflation is avoided.

Unfortunately, there is generally a tension between the above mentioned well-motivated, natural, and otherwise successful hybrid inflationary models and the recent three-year results [21] from the Wilkinson microwave anisotropy probe satellite (WMAP3). Indeed, these models, with the exception of the smooth [1] and especially the new smooth [14] hybrid inflation model, predict that,

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in global SUSY, the spectral index n_s is very close to unity and with no much running. Moreover, inclusion of supergravity (SUGRA) corrections with canonical Kähler potential yields [22], in all cases, n_s 's which are very close to unity or even exceed it. On the other hand, fitting the WMAP3 data with the standard power-law cosmological model with cold dark matter and a cosmological constant (Λ CDM), one obtains [21] n_s 's clearly lower than unity.

One possible resolution of this inconsistency is [23, 24, 25] to use a non-minimal Kähler potential with a convenient choice of the sign of one of its terms. This generates [24, 25, 26] a negative mass term for the inflaton. Consequently, the inflationary potential acquires, in general, a local maximum and minimum. Then, as the inflaton rolls from this maximum down to smaller values, hybrid inflation of the hilltop type [23] can occur. In this case, n_s can become consistent with the WMAP3 measurements, but only at the cost of a mild tuning [26] of the initial conditions. In any case, we must make sure that the system is not trapped in the local minimum of the inflationary potential, which can easily happen for general initial conditions. In such a case, no hybrid inflation would take place. Note that, in the cases of smooth and new smooth hybrid inflation, acceptable n_s 's can be obtained [14, 25] even without the appearance of this local maximum and minimum and, thus, the related complications can be avoided.

Another possibility [27] for reducing the spectral index predicted by the hybrid inflationary models is based on the observation that, in these models, n_s generally decreases with the number of e-foldings suffered by our present horizon scale during hybrid inflation. So, restricting this number of e-foldings, we can achieve values of n_s which are compatible with the recent WMAP3 data even with minimal Kähler potential. The additional number of e-foldings required for solving the horizon and flatness problems of standard hot big bang cosmology can be provided by a second stage of inflation which follows hybrid inflation. In Ref. [27], this complementary inflation was taken to be of the modular type [28] realized by a string axion at an intermediate scale. Note, in passing, that a restricted number of e-foldings during hybrid inflation was previously used [29] to achieve sufficient running of the spectral index.

In this paper, we reconsider the extended SUSY PS model of Ref. [15] which solves the b -quark mass problem and can lead to new shifted [13] or new smooth [14] hybrid inflation. We restrict ourselves in the range of parameters of this model that corresponds to the latter case. As shown in Ref. [14], the relevant scalar potential possesses, in this case, a trivial classically flat direction which is stable for large values of the inflaton field. Along this direction, the PS gauge group is unbroken. For values of the inflaton field smaller than a certain critical value, this flat direction is destabilized giving its place to a classically non-flat valley of minima along which new smooth hybrid inflation can take place. The GUT gauge group G_{PS} is broken to G_{SM} in this valley.

In Ref. [14], we investigated the possibility that all the cosmological scales exit the horizon during new smooth hybrid inflation, which is, thus, responsible for the observed spectrum of primordial fluctuations. Here, we will consider an alternative possibility. As usual, the trivial flat direction acquires [3] a logarithmic slope from one-loop radiative corrections which are due to the SUSY breaking caused by the non-vanishing potential energy density on this direction. So, a version of standard hybrid inflation can easily take place as the system slowly rolls down the trivial flat direction. We will assume here that the cosmological scales exit the horizon during this standard hybrid inflation. Then, as in Ref. [27], we can easily achieve, in global SUSY, spectral indices which are comfortably compatible with the data by restricting the number of e-foldings suffered by our present horizon scale during this inflationary period. The additional number of e-foldings required for solving the horizon and flatness problems is naturally provided, in this case, by a second stage of inflation consisting of a relatively short intermediate inflationary phase, which starts as soon as the system crosses the critical point of the trivial flat direction, followed by new smooth hybrid inflation. So, the necessary complementary inflation is automatically built in the model itself and we do not have to invoke an *ad hoc* second stage of inflation as in Ref. [27]. Moreover, large reheat temperatures can, in principle, be achieved after the second stage of inflation since this stage is realized at a superheavy scale. As a consequence, baryogenesis via thermal [30] or non-thermal [31] leptogenesis may work in this case in contrast to the model of Ref. [27], where the reheat temperature is too low for the non-perturbative electroweak sphalerons to operate. However, we should keep in mind that, as in all SUSY theories, the presence of flat directions (see e.g. Ref. [32]) can [33] naturally delay the reheating and thermalization process which follows the decay of the inflaton field. This reduces the reheat temperature and, thus, may severely constrain thermal leptogenesis. Nevertheless, non-thermal leptogenesis remains a viable alternative. Finally, the PS monopoles which are formed at the end of the standard hybrid stage of inflation can be adequately diluted by the subsequent second stage of inflation.

The inclusion of SUGRA corrections with minimal Kähler potential raises the spectral index, which, however, remains acceptable for a wide range of the model parameters. So, in this model, there is no need to include non-minimal terms in the Kähler potential and, consequently, complications from the possible appearance of a local maximum and minimum of the inflationary potential are avoided.

In Sec. II, we sketch the salient features of the extended SUSY PS model and show that it can easily lead to a successful two-stage inflationary scenario. In Sec. III, we show that this scenario remains viable even if SUGRA corrections with a minimal Kähler potential are included and, in Sec. IV, we discuss briefly gauge unification. Finally, in Sec. V, we summarize our conclusions.

II. STANDARD-SMOOTH HYBRID INFLATION IN GLOBAL SUPERSYMMETRY

We consider the extended SUSY PS model of Ref. [15]. As mentioned, this model admits a moderate violation of the asymptotic Yukawa unification so that, for $\mu > 0$, an acceptable b -quark mass is obtained even with universal boundary conditions. The breaking of G_{PS} to G_{SM} is achieved by the VEVs of the right handed neutrino type components of a conjugate pair of Higgs superfields H^c and \bar{H}^c belonging to the $(\bar{4}, 1, 2)$ and $(4, 1, 2)$ representations of G_{PS} respectively. The model also contains a gauge singlet S and a conjugate pair of superfields $\phi, \bar{\phi}$ belonging to the $(15, 1, 3)$ representation of G_{PS} . Note, in passing, that almost all SUSY inflationary models involve an *ad hoc* gauge singlet superfield such as S (an exception is the class of models of Ref. [34]). The superfield ϕ acquires a VEV which breaks G_{PS} to $G_{\text{SM}} \times \text{U}(1)_{B-L}$. In addition to G_{PS} , the model possesses a Z_2 matter parity symmetry and two global $\text{U}(1)$ symmetries, which can effectively arise [35] from the rich discrete symmetry groups encountered in many compactified string theories (see e.g. Ref. [36]). For details on the full field content and superpotential, the charge assignments, and the phenomenological and cosmological properties of this model, the reader is referred to Refs. [10, 15] (see also Ref. [16]).

The superpotential terms which are relevant for inflation are [14]

$$W = \kappa S(M^2 - \phi^2) - \gamma S H^c \bar{H}^c + m \phi \bar{\phi} - \lambda \bar{\phi} H^c \bar{H}^c, \quad (1)$$

where M, m are superheavy masses of the order of the SUSY GUT scale $M_{\text{GUT}} \simeq 2.86 \cdot 10^{16}$ GeV and κ, γ, λ are dimensionless coupling constants. All these parameters are normalized so that they correspond to the couplings between the SM singlet components of the superfields. The mass parameters M, m and any two of the three dimensionless parameters κ, γ, λ are made real and positive by appropriately redefining the phases of the superfields. The third dimensionless parameter, however, remains in general complex. For definiteness, we choose this parameter to be real and positive too as we did in Ref. [14].

The F-term scalar potential obtained from the superpotential W in Eq. (1) is given [14] by

$$V = |\kappa(M^2 - \phi^2) - \gamma H^c \bar{H}^c|^2 + |m\bar{\phi} - 2\kappa S\phi|^2 + |m\phi - \lambda H^c \bar{H}^c|^2 + |\gamma S + \lambda \bar{\phi}|^2 (|H^c|^2 + |\bar{H}^c|^2), \quad (2)$$

where the complex scalar fields which belong to the SM singlet components of the superfields are denoted by the same symbol. In Ref. [14], it was shown that this potential leads to a new version of smooth hybrid inflation provided that

$$\tilde{\mu}^2 \equiv -M^2 + \frac{m^2}{2\kappa^2} > 0 \quad (3)$$

and the parameter γ is adequately small. It was argued that, under these circumstances, there exists a trivial

classically flat direction at $\phi = \bar{\phi} = H^c = \bar{H}^c = 0$ with $V = V_{\text{tr}}^0 \equiv \kappa^2 M^4$, which is a valley of local minima for

$$|S| > S_c \equiv \sqrt{\frac{\kappa}{\gamma}} M \quad (4)$$

and becomes unstable for $|S| < S_c$, giving its place to a classically non-flat valley of minima along which new smooth hybrid inflation can take place.

We will now briefly summarize some of the main results given in Ref. [14]. The SUSY vacua of the potential in Eq. (2) lie at

$$\phi = \frac{\gamma m}{2\kappa\lambda} \left(-1 \pm \sqrt{1 + \frac{4\kappa^2\lambda^2 M^2}{\gamma^2 m^2}} \right) \equiv \phi_{\pm}, \quad (5)$$

$$\bar{\phi} = S = 0, \quad H^c \bar{H}^c = \frac{m}{\lambda} \phi. \quad (6)$$

The vanishing of the D-terms yields $\bar{H}^{c*} = e^{i\theta} H^c$, which implies that there exist two distinct continua of SUSY vacua:

$$\phi = \phi_+, \quad \bar{H}^{c*} = H^c, \quad |H^c| = \sqrt{\frac{m\phi_+}{\lambda}} \quad (\theta = 0), \quad (7)$$

$$\phi = \phi_-, \quad \bar{H}^{c*} = -H^c, \quad |H^c| = \sqrt{\frac{-m\phi_-}{\lambda}} \quad (\theta = \pi) \quad (8)$$

with $\bar{\phi} = S = 0$. One can show that the potential, besides the trivial flat direction, possesses generally two non-trivial flat directions too. One of them exists only if $\tilde{\mu}^2 < 0$ and lies at

$$\phi = \pm \sqrt{-\tilde{\mu}^2}, \quad \bar{\phi} = \frac{2\kappa\phi}{m} S, \quad H^c = \bar{H}^c = 0. \quad (9)$$

It is a shifted flat direction with $V = \kappa^2(M^4 - \tilde{\mu}^4)$ along which G_{PS} is broken to $G_{\text{SM}} \times \text{U}(1)_{B-L}$. The second non-trivial flat direction, which appears at

$$\phi = -\frac{\gamma m}{2\kappa\lambda}, \quad \bar{\phi} = -\frac{\gamma}{\lambda} S, \quad (10)$$

$$H^c \bar{H}^c = \frac{\kappa\gamma(M^2 - \phi^2) + \lambda m\phi}{\gamma^2 + \lambda^2}, \quad (11)$$

$$V = V_{\text{nsh}}^0 \equiv \frac{\kappa^2\lambda^2}{\gamma^2 + \lambda^2} \left(M^2 + \frac{\gamma^2 m^2}{4\kappa^2\lambda^2} \right)^2, \quad (12)$$

exists only for $\gamma \neq 0$ and is analogous to the trajectory for the new shifted hybrid inflation of Ref. [13]. Along this direction, G_{PS} is broken to G_{SM} . In our subsequent discussion, we will concentrate on the case $\tilde{\mu}^2 > 0$, where the shifted flat direction in Eq. (9) does not exist. It is interesting to point out that, in this case, we always have $V_{\text{nsh}}^0 > V_{\text{tr}}^0$ and it is, thus, more likely that the system will eventually settle down on the trivial rather than the new shifted flat direction.

If we expand the complex scalar fields $\phi, \bar{\phi}, H^c, \bar{H}^c$ in real and imaginary parts according to the prescription

$s = (s_1 + i s_2)/\sqrt{2}$, we find that, on the trivial flat direction, the mass-squared matrices $M_{\phi_1}^2$ of ϕ_1 , $\bar{\phi}_1$ and $M_{\phi_2}^2$ of ϕ_2 , $\bar{\phi}_2$ are

$$M_{\phi_1(\phi_2)}^2 = \begin{pmatrix} m^2 + 4\kappa^2|S|^2 \mp 2\kappa^2 M^2 & -2\kappa m S \\ -2\kappa m S & m^2 \end{pmatrix} \quad (13)$$

and the mass-squared matrices $M_{H_1}^2$ of H_1^c , \bar{H}_1^c and $M_{H_2}^2$ of H_2^c , \bar{H}_2^c are

$$M_{H_1(H_2)}^2 = \begin{pmatrix} \gamma^2|S|^2 & \mp \gamma \kappa M^2 \\ \mp \gamma \kappa M^2 & \gamma^2|S|^2 \end{pmatrix}. \quad (14)$$

The matrices $M_{\phi_1(\phi_2)}^2$ are always positive definite, while the matrices $M_{H_1(H_2)}^2$ acquire one negative eigenvalue for $|S| < S_c$. Thus, the trivial flat direction is stable for $|S| > S_c$ and unstable for $|S| < S_c$.

It has been shown in Ref. [14] that, for small enough values of the parameter γ , the trivial flat direction, after its destabilization at the critical point, gives its place to a valley of *absolute* minima for fixed $|S|$ which correspond to $\theta \simeq 0$ and lead to the SUSY vacua in Eq. (7). This valley possesses an inclination already at the classical level and can accommodate a stage of inflation with the properties of smooth hybrid inflation. In Ref. [14], the name new smooth hybrid inflation was coined for the inflationary scenario obtained when all the e-foldings required for solving the horizon and flatness problems of standard hot big bang cosmology are obtained when the system follows this valley. In this paper, we will study the case when the total required number of e-foldings splits between two stages of inflation, the standard hybrid inflation stage for $|S| > S_c$ and the new smooth hybrid inflation stage including an intermediate inflationary period for $|S| < S_c$.

The general outline of this scenario, which we call standard-smooth hybrid inflation, goes as follows. We assume that the system, possibly after a period of pre-inflation at the Planck scale, settles down on a point of the trivial flat direction with $|S| > S_c$ (see e.g. Ref. [37]). The constant classical potential energy density on this direction breaks SUSY explicitly and implies the existence of one-loop radiative corrections which lift the flatness of the potential producing the necessary inclination for driving the inflaton towards the critical point at $|S| = S_c$. So the standard hybrid inflation stage of the scenario can be realized along this path. As the system moves below the critical point, some of the masses squared of the fields become negative, resulting to a phase of spinodal decomposition. This phase is relatively fast, causes the spontaneous breaking of G_{PS} to G_{SM} , and generates a limited number of e-foldings. After this intermediate inflationary phase, the system settles down on the new smooth hybrid inflationary path and, thus, new smooth hybrid inflation takes place. The second stage of inflation consisting of the intermediate phase and the subsequent new smooth hybrid inflation yields the additional number of e-foldings required for solving the horizon and flatness problems of standard hot big bang cosmology. At the

end of this stage, the system falls rapidly into the appropriate SUSY vacuum of the theory leading, though, to no topological defect production, since the GUT gauge group is already broken to the SM gauge group during this inflationary stage. Two more requirements need to be fulfilled in order for this scenario to be viable. First, one has to make sure that the number of e-foldings generated during the second stage of inflation is adequate for diluting any monopoles generated during the phase transition at the end of the first stage of inflation. Secondly, one must ensure that all the cosmologically relevant scales receive inflationary perturbations only from the first stage of inflation so that the existence of measurable perturbations originating from the phase of spinodal decomposition, which are of a rather obscure nature, is avoided. Both of these requirements are very easily satisfied in our model, as we will see in the course of the subsequent discussion.

The one-loop radiative correction to the potential due to the SUSY breaking on the trivial inflationary path is calculated by the Coleman-Weinberg formula [38]:

$$\Delta V = \frac{1}{64\pi^2} \sum_i (-1)^{F_i} M_i^4 \ln \frac{M_i^2}{\Lambda^2}, \quad (15)$$

where the sum extends over all helicity states i , F_i and M_i^2 are the fermion number and mass squared of the i th state and Λ is a renormalization mass scale. In order to use this formula for creating a logarithmic slope in the inflationary potential, one has first to derive the mass spectrum of the model on the trivial inflationary path. It is easy to see that, in the bosonic sector, one obtains two groups of 45 pairs of real scalars with the mass-squared matrices

$$M_{-(+)}^2 = \begin{pmatrix} m^2 + 4\kappa^2|S|^2 \mp 2\kappa^2 M^2 & -2\kappa m S \\ -2\kappa m S & m^2 \end{pmatrix} \quad (16)$$

and two more groups of 8 pairs of real scalars with mass-squared matrices

$$M_{1(2)}^2 = \begin{pmatrix} \gamma^2|S|^2 & \mp \gamma \kappa M^2 \\ \mp \gamma \kappa M^2 & \gamma^2|S|^2 \end{pmatrix}. \quad (17)$$

Note that $M_{-(+)}^2$ equals $M_{\phi_1(\phi_2)}^2$ of Eq. (13) and $M_{1(2)}^2$ equals $M_{H_1(H_2)}^2$ of Eq. (14). In the fermionic sector of the theory, we obtain 45 pairs of Weyl fermions with mass-squared matrix

$$M_0^2 = \begin{pmatrix} m^2 + 4\kappa^2|S|^2 & -2\kappa m S \\ -2\kappa m S & m^2 \end{pmatrix} \quad (18)$$

and 8 more pairs of Weyl fermions with mass-squared matrix

$$\bar{M}_0^2 = \begin{pmatrix} \gamma^2|S|^2 & 0 \\ 0 & \gamma^2|S|^2 \end{pmatrix}. \quad (19)$$

The matrices M_0^2 , \bar{M}_0^2 equal $M_{-(+)}^2$, $M_{1(2)}^2$ respectively without the \mp terms in the latter matrices. The one-loop

radiative correction to the inflationary potential then takes the form

$$\begin{aligned} \Delta V = & \frac{45}{64\pi^2} \text{tr} \left\{ M_+^4 \ln \frac{M_+^2}{\Lambda^2} + M_-^4 \ln \frac{M_-^2}{\Lambda^2} \right. \\ & \left. - 2M_0^4 \ln \frac{M_0^2}{\Lambda^2} \right\} + \frac{8}{64\pi^2} \text{tr} \left\{ M_1^4 \ln \frac{M_1^2}{\Lambda^2} \right. \\ & \left. + M_2^4 \ln \frac{M_2^2}{\Lambda^2} - 2\bar{M}_0^4 \ln \frac{\bar{M}_0^2}{\Lambda^2} \right\}. \end{aligned} \quad (20)$$

The total effective potential on the trivial inflationary path will be given by $V_{\text{tr}} = v_0^4 + \Delta V$, where $v_0 \equiv \sqrt{\kappa}M$ is the inflationary scale. As already mentioned, the one-loop radiative correction to the inflationary potential lifts its classical flatness and generates a logarithmic slope which is necessary for driving the system towards the critical point at $|S| = S_c$. It is important to note that the $\sum_i (-1)^{F_i} M_i^4 = 8v_0^4 (45\kappa^2 + 4\gamma^2)$ is S -independent, which implies that the slope is Λ -independent and the scale Λ , which remains undetermined, does not enter the inflationary observables.

Making the complex scalar field S real by an appropriate global $U(1)$ R transformation and defining the canonically normalized real inflaton field $\sigma \equiv \sqrt{2}S$, the slow-roll parameters ε , η and the parameter ξ^2 , which enters the running of the spectral index, are (see e.g. Ref. [39])

$$\varepsilon \equiv \frac{m_{\text{P}}^2}{2} \left(\frac{V'(\sigma)}{V(\sigma)} \right)^2, \quad (21)$$

$$\eta \equiv m_{\text{P}}^2 \left(\frac{V''(\sigma)}{V(\sigma)} \right), \quad (22)$$

$$\xi^2 \equiv m_{\text{P}}^4 \left(\frac{V'(\sigma)V'''(\sigma)}{V^2(\sigma)} \right), \quad (23)$$

where the prime denotes derivation with respect to the inflaton σ and $m_{\text{P}} \simeq 2.44 \cdot 10^{18}$ GeV is the reduced Planck mass. In these equations, V is either the effective potential V_{tr} on the trivial inflationary path defined above, if we are referring to the standard hybrid stage of inflation, or the effective potential V_{nsm} for new smooth hybrid inflation, which has to be calculated numerically (see Ref. [14]), if we are referring to the new smooth hybrid inflationary phase.

Numerical simulations have shown that, after crossing the critical point at $\sigma = \sigma_c \equiv \sqrt{2}S_c$, the system continues evolving, for a while, with the Hubble parameter H remaining approximately constant and equal to $H_0 \equiv v_0^2/\sqrt{3}m_{\text{P}}$ until it settles down on the new smooth hybrid inflationary path at $\sigma \approx 0.99\sigma_c$. The scale factor of the universe increases by about 8 e-foldings during this intermediate period. The fields H^c and \bar{H}^c are effectively massless at $\sigma = \sigma_c$ and, thus, acquire inflationary perturbations $\delta H^c = \delta \bar{H}^c \approx H_0/2\pi$. Their initial values at the critical point are taken equal to these perturbations. The inflaton σ is assumed to have an initial velocity given by the slow-roll equation

$$\dot{\sigma} = -\frac{V'_{\text{tr}}(\sigma_c)}{3H_0}, \quad (24)$$

where the overdot denotes derivation with respect to the cosmic time t and the inclination $V'_{\text{tr}}(\sigma_c)$ is provided by the radiative corrections on the trivial flat direction (for the parameter values that are of interest, the slow-roll conditions $\varepsilon \leq 1$, $|\eta| \leq 1$ for the first stage of inflation are violated only “infinitesimally” close to the critical point). Although the above results are not independent from the values of the model parameters, they represent legitimate mean values. Moreover, inflationary observables like the spectral index have shown not to depend significantly on the properties of this intermediate phase.

From the above discussion, we see that the number of e-foldings from the time when the pivot scale $k_0 = 0.002 \text{ Mpc}^{-1}$ crosses outside the inflationary horizon until the end of inflation is (see e.g. Ref. [39])

$$\begin{aligned} N_Q \approx & \frac{1}{m_{\text{P}}^2} \int_{\sigma_f}^{0.99\sigma_c} \frac{V_{\text{nsm}}(\sigma)}{V'_{\text{nsm}}(\sigma)} d\sigma + 8 \\ & + \frac{1}{m_{\text{P}}^2} \int_{\sigma_c}^{\sigma_Q} \frac{V_{\text{tr}}(\sigma)}{V'_{\text{tr}}(\sigma)} d\sigma, \end{aligned} \quad (25)$$

where $\sigma_Q \equiv \sqrt{2}S_Q > 0$ is the value of the inflaton field at horizon crossing of the pivot scale and σ_f refers to the value of σ at the end of the second stage of inflation and can be found from the corresponding slow-roll conditions. The power spectrum $P_{\mathcal{R}}$ of the primordial curvature perturbation at the scale k_0 is given (see e.g. Ref. [39]) by

$$P_{\mathcal{R}}^{1/2} \simeq \frac{1}{2\pi\sqrt{3}} \frac{V_{\text{tr}}^{3/2}(\sigma_Q)}{m_{\text{P}}^3 V'_{\text{tr}}(\sigma_Q)}. \quad (26)$$

The spectral index n_s , the tensor-to-scalar ratio r , and the running of the spectral index $dn_s/d\ln k$ can be written (see e.g. Ref. [39]) as

$$\begin{aligned} n_s \simeq & 1 + 2\eta - 6\varepsilon, \quad r \simeq 16\varepsilon, \\ \frac{dn_s}{d\ln k} \simeq & 16\varepsilon\eta - 24\varepsilon^2 - 2\xi^2, \end{aligned} \quad (27)$$

where ε , η , and ξ^2 are evaluated at $\sigma = \sigma_Q$. The number of e-foldings N_Q required for solving the horizon and flatness problems of standard hot big bang cosmology is given (see e.g. Ref. [40]) approximately by

$$N_Q \simeq 53.76 + \frac{2}{3} \ln \left(\frac{v_0}{10^{15} \text{ GeV}} \right) + \frac{1}{3} \ln \left(\frac{T_{\text{r}}}{10^9 \text{ GeV}} \right), \quad (28)$$

where T_{r} is the reheat temperature that is expected not to exceed about 10^9 GeV, which is the well-known gravitino bound [41].

As already explained, magnetic monopoles are produced at the end of the standard hybrid stage of inflation, where G_{PS} breaks down to G_{SM} . We will now discuss, in some detail, this production of magnetic monopoles and their dilution by the subsequent second stage of inflation. The masses of the fields H^c and \bar{H}^c , which vanish at $\sigma = \sigma_c$, grow very fast as the system moves to smaller

values of σ . Actually, as one can show numerically, they become of order H_0 when the system is still “infinitesimally” close to the critical point and the inflationary perturbations of H^c and \bar{H}^c become suppressed. After this, the system evolves essentially classically. It remains, for a while, close to the trivial flat direction (which, for $\sigma < \sigma_c$, is unstable as it consists of saddle points) yielding about 8 e-foldings as mentioned above. It, finally, settles down on the new smooth hybrid inflationary path at $\sigma \approx 0.99 \sigma_c$. To be more precise, it ends up at a point of the manifold which consists of the *absolute* minima of the potential for fixed $\sigma \approx 0.99 \sigma_c$. The particular choice of this point is made by the inflationary perturbations of H^c and \bar{H}^c , which cease to operate when the masses of these fields reach the value H_0 . This happens after crossing the critical point, but “infinitesimally” close to it, as we already mentioned. So the correlation length which is relevant for magnetic monopole production by the Kibble mechanism [42] is $\approx H_0^{-1}$.

The initial monopole number density can then be estimated [42] as

$$n_M^{\text{init}} \approx \frac{3\mathbf{p}}{4\pi} H_0^3, \quad (29)$$

where $\mathbf{p} \sim 1/10$ is a geometric factor. At the end of inflation, the monopole number density becomes

$$n_M^{\text{fin}} \approx \frac{3\mathbf{p}}{4\pi} H_0^3 e^{-3\delta N}, \quad (30)$$

where δN is the total number of e-foldings during the intermediate period and the subsequent new smooth hybrid inflation phase. Dividing n_M^{fin} by the number density $n_{\text{infl}} \approx V_{\text{tr}}^0/m_{\text{infl}}$ of the inflatons which are produced at the termination of inflation (m_{infl} is the inflaton mass), we obtain that, at the end of inflation, the number density of monopoles n_M is given by

$$\frac{n_M}{n_{\text{infl}}} \approx \frac{3\mathbf{p}}{4\pi} H_0^3 e^{-3\delta N} \frac{m_{\text{infl}}}{V_{\text{tr}}^0}. \quad (31)$$

This ratio remains practically constant until reheating, where the relative number density of monopoles can be estimated as (compare with Ref. [43])

$$\frac{n_M}{s} = \frac{n_M}{n_{\text{infl}}} \frac{n_{\text{infl}}}{s} \approx \frac{3\mathbf{p}}{16\pi} \frac{H_0 T_r}{m_P^2} e^{-3\delta N}, \quad (32)$$

where s is the entropy density and the relations $n_{\text{infl}}/s = 3T_r/4m_{\text{infl}}$ (in the instantaneous inflaton decay approximation) and $3H_0^2 = V_{\text{tr}}^0/m_P^2$ were used. After reheating, the relative number density of monopoles remains essentially unaltered provided that there is no entropy production at subsequent times. Taking $n_M/s \lesssim 10^{-30}$, which corresponds [44] to the Parker bound [45] on the present magnetic monopole flux in our galaxy derived from galactic magnetic field considerations, $T_r \sim 10^9 \text{ GeV}$, and $H_0 \sim 10^{12} \text{ GeV}$, we obtain from Eq. (32) that $\delta N \gtrsim 9.2$. Using Eq. (28), this implies that $N_{\text{st}} \lesssim 45$, where N_{st}

is the number of e-foldings of the pivot scale k_0 during the standard hybrid stage of inflation. Saturating this bound, we obtain a monopole flux which may be measurable. However, the interesting values of N_{st} encountered here in the global SUSY case are much smaller (see below) and, thus, the predicted magnetic monopole flux is unlikely to be measurable. In the minimal SUGRA case, N_{st} is restricted to quite small values (see Sec. III) and the monopole flux is predicted utterly negligible.

The model contains five free parameters, namely M , m , κ , γ , and λ . As already mentioned, the VEVs of H^c , \bar{H}^c break the PS gauge group to G_{SM} , whereas the VEV of the field ϕ breaks it only to $G_{\text{SM}} \times \text{U}(1)_{B-L}$. So, the gauge boson A^\perp corresponding to the linear combination of $\text{U}(1)_Y$ and $\text{U}(1)_{B-L}$ which is perpendicular to $\text{U}(1)_Y$ acquires its mass squared $m_{A^\perp}^2 = (5/2)g^2|\langle H^c \rangle|^2$ solely from the VEVs $\langle H^c \rangle$, $\langle \bar{H}^c \rangle$ of H^c , \bar{H}^c (g is the SUSY GUT gauge coupling constant). On the other hand, the masses squared m_A^2 and $m_{W_R}^2$ of the color triplet, anti-triplet (A^\pm) and charged $\text{SU}(2)_R$ (W_R^\pm) gauge bosons get contributions from $\langle \phi \rangle$ too. Namely, $m_A^2 = g^2(|\langle H^c \rangle|^2 + (4/3)|\langle \phi \rangle|^2)$ and $m_{W_R}^2 = g^2(|\langle H^c \rangle|^2 + 2|\langle \phi \rangle|^2)$. As we will see below, the VEVs of H^c and ϕ in the SUSY vacua of the model turn out to be of the same order of magnitude. Since the A^\pm gauge bosons are expected to affect the renormalization group equations to a greater extent than the W_R^\pm ones (the SM singlet gauge boson A^\perp does not affect them at all), we set the mass m_A divided by $g \approx 0.7$ equal to the SUSY GUT scale M_{GUT} . We also set the value of the parameter $p \equiv \sqrt{2}\kappa M/m$ equal to $1/\sqrt{2}$. Note that, for $\tilde{\mu}^2 > 0$, this parameter is smaller than unity as seen from Eq. (3). Finally, we take T_r to saturate the gravitino bound [41], i.e. $T_r \simeq 10^9 \text{ GeV}$, and fix the power spectrum of the primordial curvature perturbation to the WMAP3 [21] normalization $P_{\mathcal{R}}^{1/2} \simeq 4.85 \cdot 10^{-5}$ at the pivot scale k_0 . These choices fix three of the five parameters of the model. So, we are left with two free parameters. We will take the ratio $\alpha \equiv |\langle H^c \rangle|/|\langle \phi \rangle|$, which, for γ adequately small, approximately equals $\sqrt{m/\lambda M}$, to be one of them. The second free parameter can be chosen to be the number of e-foldings N_{st} of the pivot scale k_0 during the standard hybrid stage of inflation (N_{st} can be fixed by adjusting e.g. the parameter γ). We will plot our results as functions of these two free parameters.

In Fig. 1, we plot the predicted spectral index of the model versus the number of e-foldings N_{st} suffered by the pivot scale k_0 during the standard hybrid stage of inflation for various values of the parameter α . Note that N_{st} is given by the last term in the right-hand side of Eq. (25). We have restricted ourselves to N_{st} 's between 4 and 45. The lower limit guarantees the validity of our requirement that all the cosmological scales receive perturbations from the first stage of inflation. Indeed, the number of e-foldings that elapse between the horizon crossing of the pivot scale k_0 and the largest cosmological scale $0.1/\text{Mpc}$ is about 4. The upper limit on N_{st} ensures that the present flux of magnetic monopoles in our galaxy does not exceed the Parker bound [45] as we

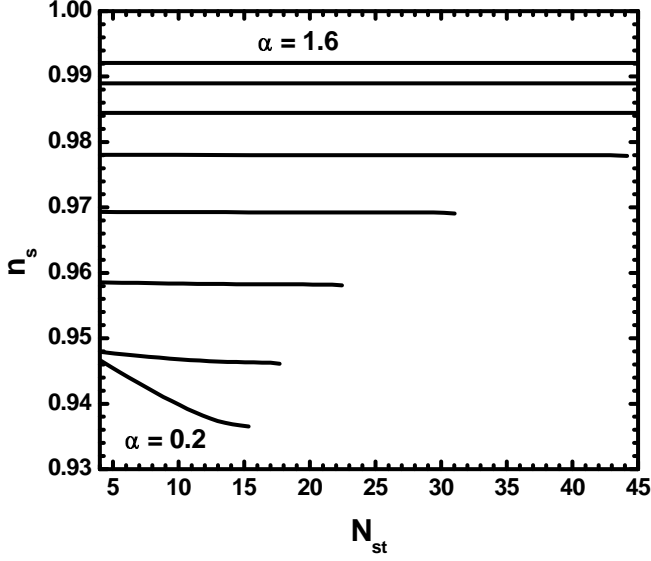


FIG. 1: Spectral index in standard-smooth hybrid inflation versus N_{st} in global SUSY for $p = \sqrt{2}\kappa M/m = 1/\sqrt{2}$. The values of the parameter $\alpha = \langle H^c \rangle / |\langle \phi \rangle|$ range from 0.2 to 1.6 with steps of 0.2.

showed above. The parameter α is limited between 0.2 and 1.6. Values of α lower than about 0.2 require non-perturbative values of λ , whereas $\alpha = 1.6$ or higher is of no much interest since it leads to unacceptably large n_s 's. Whenever a curve in Fig. 1 terminates on the right, this means that the constraint $P_{\mathcal{R}}^{1/2} \simeq 4.85 \cdot 10^{-5}$ cannot be satisfied beyond this endpoint. The WMAP3 data fitted by the standard power-law Λ CDM cosmological model predict [21] that, at the pivot scale k_0 ,

$$n_s = 0.958 \pm 0.016 \Rightarrow 0.926 \lesssim n_s \lesssim 0.99 \quad (33)$$

at 95% confidence level. We see, from Fig. 1, that one can readily obtain from our model spectral indices which lie within this 2- σ allowed range. Moreover, the 1- σ range is fully covered by the predicted values of n_s . Note, however, that one cannot obtain spectral indices lower than about 0.936. It is obvious that large values of N_{st} are of no much interest since they yield large n_s 's. So a possibly measurable flux of monopoles at the level of the Parker bound is very unlikely.

For the curves depicted in Fig. 1, γ varies in the range $\gamma \simeq (0.04 - 6) \cdot 10^{-3}$. It increases as α decreases or N_{st} increases with its dependence on N_{st} being much milder. The ranges of the other parameters of the model are $\kappa \simeq (0.46 - 3.62) \cdot 10^{-2}$, $\lambda \simeq 0.004 - 1.56$, $M \simeq (1.45 - 2.44) \cdot 10^{16}$ GeV, $m \simeq (0.13 - 1.56) \cdot 10^{15}$ GeV, $\sigma_Q \simeq (0.9 - 8.8) \cdot 10^{17}$ GeV, $\sigma_c \simeq (0.8 - 2.3) \cdot 10^{17}$ GeV, and $\sigma_f \simeq (0.5 - 1.5) \cdot 10^{17}$ GeV. The total number of e-foldings from the time when the pivot scale k_0 crosses outside the inflationary horizon until the end of the second stage of inflation is $N_Q \simeq 53.7 - 54.7$. Finally, $dn_s/d \ln k \simeq -(0.06 - 4) \cdot 10^{-3}$ and the tensor-to-scalar ratio $r \simeq (0.008 - 2.8) \cdot 10^{-4}$. A decrease in the value

of p , which is the only arbitrarily chosen parameter, generally leads to an increase of the spectral index. Thus, smaller values of p are expected to shift the curves in Fig. 1 upwards, but otherwise do not change the qualitative features of the model.

III. SUPERGRAVITY CORRECTIONS

We now turn to the discussion of the SUGRA corrections to the inflationary potentials of our model. The F-term scalar potential in SUGRA is given by

$$V = e^{K/m_{\text{P}}^2} \left[(F_i)^* K^{i^*j} F_j - 3 \frac{|W|^2}{m_{\text{P}}^2} \right], \quad (34)$$

where K is the Kähler potential, $F_i = W_i + K_i W/m_{\text{P}}^2$, a subscript i (i^*) denotes derivation with respect to the complex scalar field s^i (s^{i^*}) and K^{i^*j} is the inverse of the Kähler metric K_{ji^*} . We will only consider SUGRA with minimal Kähler potential and show that the WMAP3 results [21] can be met for a wide range of values of the parameters of the model.

The minimal Kähler potential in the model under consideration has the form

$$K^{\text{min}} = |S|^2 + |\phi|^2 + |\bar{\phi}|^2 + |H^c|^2 + |\bar{H}^c|^2 \quad (35)$$

and the corresponding F-term scalar potential is

$$V^{\text{min}} = e^{K^{\text{min}}/m_{\text{P}}^2} \left[\sum_s \left| W_s + \frac{W s^*}{m_{\text{P}}^2} \right|^2 - 3 \frac{|W|^2}{m_{\text{P}}^2} \right], \quad (36)$$

where s stands for any of the five complex scalar fields appearing in Eq. (35). It is very easily verified that, on the trivial flat direction, this scalar potential expanded up to fourth order in $|S|$ takes the form

$$V_{\text{tr}}^{\text{min}} \simeq v_0^4 \left(1 + \frac{1}{2} \frac{|S|^4}{m_{\text{P}}^4} \right). \quad (37)$$

Thus, after including the SUGRA corrections with minimal Kähler potential, the effective potential during the standard hybrid stage of inflation becomes

$$V_{\text{tr}}^{\text{SUGRA}} \simeq V_{\text{tr}}^{\text{min}} + \Delta V \quad (38)$$

with ΔV representing the one-loop radiative correction given in Eq. (20). Furthermore, it has been shown in Ref. [14] that the effective potential on the new smooth hybrid inflationary path in the presence of minimal SUGRA takes the form

$$V_{\text{nsm}}^{\text{SUGRA}} \simeq v_0^4 \left(\tilde{V}_{\text{nsm}} + \frac{1}{2} \frac{|S|^4}{m_{\text{P}}^4} \right), \quad (39)$$

where $\tilde{V}_{\text{nsm}} \equiv V_{\text{nsm}}/v_0^4$ with V_{nsm} being the effective potential on the new smooth hybrid inflationary path in the case of global SUSY. Note that, in the minimal SUGRA

case, the critical value of σ , where the trivial flat direction becomes unstable, will be slightly different from the critical value of σ in the global SUSY case.

The cosmology of the model after including the minimal SUGRA corrections follows straightforwardly from that of the global SUSY case if one replaces the inflationary effective potentials of the latter by the ones derived above and take into account some changes in the intermediate phase between the two main inflationary periods. Actually, one finds numerically that, due to the larger inclination of the inflationary path provided by the minimal SUGRA corrections, the number of e-foldings during the intermediate period of inflation is reduced to about 2 or 3. Also, the value of σ at which the system settles down on the new smooth hybrid inflationary path decreases to about $\sigma \approx 0.95 \sigma_c$. Moreover, as it turns out, the evolution of the system can be very well approximated by the simplifying assumption that, during the intermediate phase, it follows the new smooth hybrid inflationary path. Therefore, we remove the term 8 from the right-hand side of Eq. (25) and replace the upper limit in the first integral by σ_c .

We again set the mass m_A of the color triplet, anti-triplet gauge bosons divided by $g \approx 0.7$ equal to the SUSY GUT scale M_{GUT} and the value of the parameter $p = \sqrt{2}\kappa M/m$ equal to $1/\sqrt{2}$. We also take T_r to saturate the gravitino bound [41], i.e. $T_r \simeq 10^9$ GeV, and fix the power spectrum of the primordial curvature perturbation to the WMAP3 [21] normalization $P_{\mathcal{R}}^{1/2} \simeq 4.85 \cdot 10^{-5}$ at the pivot scale k_0 . Finally, we will again plot our results against the parameter $\alpha = |\langle H^c \rangle|/|\langle \phi \rangle|$ and the number of e-foldings N_{st} of the pivot scale k_0 during the standard hybrid stage of inflation (N_{st} can again be fixed by adjusting e.g. the parameter γ).

In Fig. 2, we plot the predicted spectral index of the model in minimal SUGRA versus N_{st} for various values of the parameter α . We have allowed N_{st} to vary only between 4 and 45 for the same reasons mentioned in the global SUSY case. For α smaller than about 0.2, the required values of λ turn out again to be non-perturbative, whereas, for α greater than about 0.7, the WMAP3 normalization of the power spectrum of the primordial curvature perturbation is not satisfied. We see that spectral indices below unity are readily obtainable and that the central value $n_s = 0.958$ from the WMAP3 results is achievable. Though, the spectral index cannot be reduced below $n_s \simeq 0.953$, as is evident from the curve with $\alpha = 0.2$. Note that values of n_s in the 95% confidence level range of Eq. (33) can be obtained only if N_{st} is lower than about 21. So, the predicted magnetic monopole flux in our galaxy is utterly negligible.

The range of variance of the parameter γ on the curves of Fig. 2 is $\gamma \simeq (0.17 - 3.43) \cdot 10^{-3}$ with γ increasing with decreasing α and slightly increasing with increasing N_{st} . The ranges of the other parameters of the model on these curves are $\kappa \simeq (0.66 - 1.35) \cdot 10^{-2}$, $\lambda \simeq 0.027 - 0.68$, $M \simeq (2.12 - 2.44) \cdot 10^{16}$ GeV, $m \simeq (2.8 - 6.6) \cdot 10^{14}$ GeV, $\sigma_Q \simeq (0.95 - 3.05) \cdot 10^{17}$ GeV, $\sigma_c \simeq (0.6 - 2) \cdot 10^{17}$ GeV, and

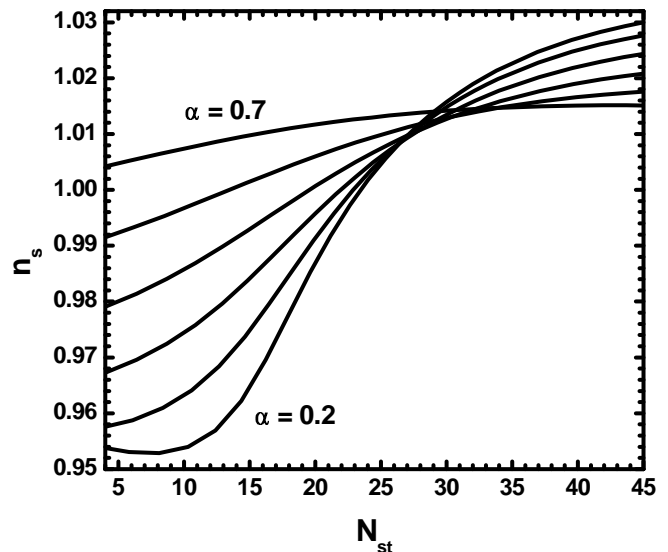


FIG. 2: Spectral index in standard-smooth hybrid inflation versus N_{st} in minimal SUGRA for $p = \sqrt{2}\kappa M/m = 1/\sqrt{2}$. The values of the parameter α range from 0.2 to 0.7 with steps of 0.1.

$\sigma_f \simeq (4.9 - 9.9) \cdot 10^{16}$ GeV. The total number of e-foldings from the time when the pivot scale k_0 crosses outside the inflationary horizon until the end of the second stage of inflation is $N_Q \simeq 54.1 - 54.5$. Finally, $dn_{rms}/d \ln k \simeq -(0.77 - 3.76) \cdot 10^{-3}$ and $r \simeq (0.7 - 5.3) \cdot 10^{-5}$. Again, a decrease in the value of p generally leads to an increase of the spectral index, resulting, thus, to a shift of the curves in Fig. 2 upwards. However, the other qualitative features of the model are not affected.

IV. GAUGE UNIFICATION

We will now briefly address the question of gauge unification in our model. As the careful reader may have noticed, cosmological considerations have constrained the mass parameter m to be significantly lower than M_{GUT} , especially in the case of minimal SUGRA. This could easily jeopardize the unification of gauge coupling constants and, indeed, it does, as it turns out, since some of the fields that contribute significantly to the gauge coupling constant running acquire masses of order m . Actually, there are two different scales below M_{GUT} that give masses to fields contributing to the renormalization group equations for the gauge coupling constants. One of them is, as already mentioned, around m and the other is around $|\langle H^c \rangle| = \sqrt{m|\langle \phi \rangle|/\lambda}$. This holds in the minimal SUGRA case and, for not too large n_s 's, in the global SUSY case too. Gauge unification is destroyed for two reasons. First of all, the fields which acquire masses below M_{GUT} are too many and this causes the appearance of Landau poles in the running of the gauge coupling constants. Secondly, none of these fields has $\text{SU}(2)_L$

quantum numbers and thus, even if divergences were not present, the $SU(2)_L$ gauge coupling constant would fail to unify with the other gauge coupling constants.

The first problem is avoided by considering the superpotential term $\xi\phi^2\bar{\phi}$, which is allowed by all the symmetries of the theory (see Ref. [15]). The reason for not including this term in our discussion from the beginning is that it does not contain a coupling between the SM singlet components of ϕ , $\bar{\phi}$ and so does not affect the inflationary dynamics. This is because $\phi^2\bar{\phi}$ is the mixed product of the three vectors ϕ , ϕ , and $\bar{\phi}$ in the 3-dimensional space in which the $SO(3)$ group which is locally equivalent to $SU(2)_R$ operates. Nevertheless, this term generates extra contributions of order $|\xi\langle\phi\rangle|^2$ to the masses squared of some fields and, thus, helps us to get rid of the Landau poles.

The second problem can be solved only by including extra fields in the model which affect the running of the $SU(2)_L$ gauge coupling constant. Note that, although the extended PS model under consideration already contains fields with $SU(2)_L$ quantum numbers which are not present in the minimal SUSY PS model, namely the fields h' and \bar{h}' belonging to the $(15,2,2)$ representation (see Ref. [15]), these fields are not sufficient for achieving the desired gauge unification since they do not affect the running of the $SU(2)_L$ gauge coupling constant as much as it is required. Consequently, one has to consider the inclusion of some extra fields. There is a good choice which uses a single extra field, namely a superfield χ belonging to the $(15,3,1)$ representation. If we require that this field has charge 1/2 under the global $U(1)_R$ symmetry, then the only superpotential term in which this field is allowed to participate is a mass term of the form $\frac{1}{2}m_\chi\chi^2$. One can then tune the new mass parameter m_χ so as to achieve unification of the gauge coupling constants. We find that this mass should be $\approx 8 \cdot 10^{14}$ GeV.

It turns out that one can achieve gauge unification at the appropriate scale ($\approx 2 \cdot 10^{16}$ GeV) as long as the mass parameter m is constrained to lie above $3 \cdot 10^{14}$ GeV. This condition is fulfilled for almost all curves of Figs. 1 and 2 except for the curves with $\alpha = 1.2, 1.4$, and 1.6 in Fig. 1. Note that this constraint is equivalent to the statement that the spectral index in the global SUSY case is less than about 0.98. So, the low spectral index regime is not affected. Furthermore, if one wants to be on the safe side avoiding marginal gauge unification (the value $m \approx 3 \cdot 10^{14}$ GeV leads to gauge unification with a rather large GUT gauge coupling constant, which is of order unity or larger), then one can impose the restriction $m \gtrsim 4 \cdot 10^{14}$ GeV, which leads to the constraints $\alpha \lesssim 0.8$ for Fig. 1 and $\alpha \lesssim 0.5$ for Fig. 2.

V. CONCLUSIONS

We have reconsidered the extended SUSY PS model of Ref. [15] which solves the b -quark mass problem. In this model, exact asymptotic Yukawa unification is nat-

urally and moderately violated so that, for $\mu > 0$, the predicted b -quark mass lies within the experimentally allowed range even with universal boundary conditions. The same model can automatically lead to new versions of the shifted and smooth hybrid inflationary scenarios based solely on renormalizable superpotential interactions. In both of these cases, the PS GUT gauge group is broken to the SM gauge group already during inflation and, thus, no PS magnetic monopole production takes place at the end of inflation. So, the possible cosmological catastrophe from magnetic monopole overproduction is avoided. In contrast to new smooth hybrid inflation, the new shifted one yields, in global SUSY, spectral indices which are too close to unity and without much running in conflict with the recent WMAP3 data. Moreover, inclusion of minimal SUGRA raises n_s to unacceptably large values in both of these inflationary scenarios.

To resolve this problem, we proposed a two-stage inflationary scenario which is naturally realized within this extended SUSY PS model for the range of its parameters leading to new smooth hybrid inflation. The first stage of inflation is of the standard hybrid type and takes place along the trivial classically flat direction of the scalar potential, which is stable for values of the inflaton field larger than a certain critical value. The inflaton is driven by the logarithmic slope acquired by this direction from one-loop radiative corrections which are due to the SUSY breaking caused by the non-vanishing potential energy density on this direction. Note that, on the trivial flat direction, the PS gauge group is unbroken. Assuming that the cosmological scales exit the horizon during the first stage of inflation, we can achieve, in global SUSY, spectral indices compatible with the WMAP3 data by restricting the number of e-foldings suffered by our present horizon scale during this inflationary stage.

The system, after crossing the critical point of the trivial flat direction, undergoes a relatively short intermediate inflationary phase and then falls rapidly into the new smooth hybrid inflationary path along which it continues inflating as it slowly rolls towards the vacua. Note that this path appears right after the destabilization of the trivial flat direction at its critical point. During this second stage of (intermediate plus new smooth hybrid) inflation, the additional number of e-foldings needed for solving the horizon and flatness problems is naturally generated and G_{PS} is broken to G_{SM} . So, we see that the necessary complementary inflation is automatically built in the model itself and we do not have to invoke an *ad hoc* second stage of inflation as in other scenarios. Moreover, large reheat temperatures can be achieved after the second stage of inflation since this stage is realized at a superheavy scale. Therefore, baryogenesis via (non-thermal) leptogenesis may work in this case in contrast to other models where the reheat temperature is too low for sphalerons to operate. Finally, the PS monopoles that are formed at the end of the standard hybrid stage of inflation can be adequately diluted by the second stage of inflation. The monopole flux in our galaxy in the case of

global SUSY is expected to be utterly negligible for not too large values of the spectral index.

Including SUGRA corrections with minimal Kähler potential enhances the predicted values of the spectral index, which, however, remain within the allowed interval for a wide range of the model parameters. So, in this model, there is no need to include non-minimal terms in the Kähler potential and, thus, complications from the possible appearance of a local maximum and minimum of the inflationary potential are avoided. The monopole flux in the SUGRA case turns out not to be measurable

for all the allowed values of the model parameters.

ACKNOWLEDGEMENTS

We thank I.N.R. Peddie for his help with gauge unification. This work was supported by the European Union under the contracts MRTN-CT-2004-503369 and HPRN-CT-2006-035863 as well as the Greek Ministry of Education and Religion and the EPEAK program Pythagoras.

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